

Experimental Challenges for Quantum Gravity¹

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ABSTRACT

The existence of a new fundamental scale may lead to modified dispersion relations for particles at high energies. Such modifications seem to be realized with the Planck scale in certain descriptions of quantum gravity. We apply effective field theory to this problem and identify dimension 5 operators that would lead to cubic modifications of dispersion relations for Standard Model particles. We also discuss other issues related to this approach including various experimental bounds on the strength of these interactions. Further we sketch a scenario where mixing of these operators with dimensions 3 and 4 due to quantum effects is minimal.

¹Talk presented by RCM at *QTS3: Third International Symposium on Quantum Theory and Symmetries*, held at the University of Cincinnati, September 10-14, 2003. To appear in the Proceedings.

1 Introduction

The Planck mass, M_{Pl} , the dimensional parameter determining the strength of gravitational interactions, remains a source of conceptual problems for quantum field theories. When the momentum transfer in two particle collisions is comparable to the Planck mass, the graviton exchange becomes strong, signifying a breakdown of the perturbative field theory description. Even without having a fully consistent fundamental theory at hand, one can hypothesize several broad categories of the low-energy effects induced by M_{Pl} . The first group of such theories has only minor modifications due to the existence of new physics at M_{Pl} . By minor modifications we understand that all “sacred” symmetries (Lorentz symmetry, CPT, spin-statistics, etc.) of the field theory remain unbroken at low energies. Critical string theory in simple backgrounds reduces to a field theory plus gravity at the scales lower than M_s and so provides an interesting example of this category. In this case the chances to probe $1/M_{\text{Pl}}$ effects are very remote as it is insufficient to merely have a probe with large energy. Rather one must have extremely large momentum transfers. Consequently, the propagation of a free particle with large energy/momentum is immune to the effects of new physics, as all corrections to the dispersion relation could be cast in the form of $(p^2)(p^{2n}/M_{\text{Pl}}^{2n}) = (m^{2n+2}/M_{\text{Pl}}^{2n})$ where m is the mass of the particle and p is the four-momentum.

The hope that nature could be more gracious to physicists is reflected in the second class of scenarios, where the $1/M_{\text{Pl}}$ effects have much more “vivid” properties. Loop quantum gravity seems to provide an example of this scenario. In this approach, the discrete nature of space at short distances may be expected to induce violations of Lorentz invariance and CPT. Such violations are also often discussed in a broader context using field theoretical language[1]. Here one might assume a perfectly Lorentz symmetric, CPT conserving action at $1/M_{\text{Pl}}^0$ order and account for the existence of the Lorentz breaking terms via a set of higher dimension operators. This would lead to modifications of the dispersion relation for a free particle as terms of the form E^{n+2}/M_{Pl}^n can appear. Such effects can be searched for both with astrophysical observations[2, 3, 4, 5, 6] and with high precision low-energy experiments[7, 8, 9].

Cubic modifications of dispersion relations, which would appear at the leading order in $1/M_{\text{Pl}}$, have received considerable attention in the literature recently[2, 3, 4, 5, 6, 10, 11]. In the language of the effective field theory, such modifications can be described by the dimension 5 operators. Although dimension 3 and 4 operators were extensively studied[1], dimension 5 operators remained unclassified and poorly explored. In our recent work[9], we took a first step towards such a classification by considering dimension 5 operators in which Lorentz breaking can be achieved by the introduction of a background four-vector n^a . Within this framework, we studied cubic modifications of the dispersion relation for scalars, vectors and fermions. Specializing the operators to Standard Model particles, our results showed that a cubic modification of the dispersion relation is not possible for the Higgs particles, must have an opposite sign for opposite chiralities of photons, and is independent for different chiralities in the case of fermions.

This paper provides a summary of our earlier results[9]. We begin with a review the class of dimension 5 operators which are of interest to produce cubic modifications of dispersion relations. We also discuss experimental constraints limiting the strength of these operators.

In particular, we demonstrate remarkably stringent bounds[9] on these operators coming from terrestrial clock comparison experiments.

Beyond this summary, we address further issues vis-a-vis Lorentz violation and its effective field theory description in the discussion section. In particular, we speculate on possible mechanisms by which the violation of Lorentz invariance may be restricted to higher dimensional operators and address the mixing of the dimension 5 operators with those of dimension 3 and 4 through quantum effects. We are able to identify a scenario in which this mixing may be minimal. We also argue that although our previous discussion is framed in the language of Lorentz symmetry breaking, many of the results still have application in scenarios where the symmetry is deformed[11]. Finally, we stress the importance of a consistent formulation of Lorentz-violating (or Lorentz-deforming) theories in the presence of gravity[12] and the necessity of correct identification of the infrared degrees of freedom in such theories.

2 Dimension 5 operators and cubic dispersion relations

In the framework of low energy effective field theory, the modified dispersion relations should be derived from some appropriate modification of the kinetic terms in the Lagrangian. Such modification may appear in the leading dimension 3 and 4 terms. However, we assume that short-distance physics does not generate such Lorentz violating operators directly. From a pragmatic point of view, this assumption is quite safe since even if these terms exist, experimental constraints indicate that they must be exceptionally small[1]. Nevertheless, it poses serious theoretical problems and we will attempt to address this point in the closing discussion.

At the next level are dimension 5 operators which would lead to $O(E^3)$ modifications of the dispersion relations. We adopt the simplest approach where Lorentz symmetry is broken by a background four-vector n^a (with $n \cdot n = 1$). We construct operators satisfying six generic criteria:

1. Quadratic in the *same* field
2. One more derivative than the usual kinetic term
3. Gauge invariant
4. Lorentz invariant, except for the appearance of n^a
5. Not reducible to lower dimension by the equations of motion
6. Not reducible to a total derivative

Conditions 2 and 5 ensure that these operators lead to $O(E^3)$ modifications of the dispersion relations, rather than $O(E^2 m)$ or $O(E m^2)$, where m is the mass of the particle. Our working assumption will also be that these operators are naturally suppressed by a factor of $1/M_{\text{Pl}}$, and that $m \ll E \ll M_{\text{Pl}}$. This scaling ensures that all operators of dimension 5 can be regarded as small perturbations. Below, we consider the cases of vector and fermion particles. We refer the interested reader to Ref. [9] for the analogous discussion of scalars.

Vector: Consider a U(1) gauge field with the leading kinetic term, $\mathcal{L}_0 = -F^2/4$. The leading order equations of motion are just the Maxwell equations, $\partial_a F^{ab} = 0$. After gauge fixing $\partial \cdot A = 0$, this yields $\square A_a = 0$ or $k^2 A_a(k) = 0$ in momentum space with $A_a \sim \exp(ik \cdot x)$. We wish to modify the dispersion relation at $O(E^3)$ and so the new terms should satisfy the constraints listed above. Keeping in mind the leading order Maxwell equations and the Bianchi identities $\partial_{[a} F_{bc]} = 0$, one finds that there is a *unique* term with the desired properties

$$\mathcal{L}_\gamma = \frac{\xi}{M_{\text{Pl}}} n^a F_{ad} n \cdot \partial (n_b \tilde{F}^{bd}), \quad (1)$$

where $\tilde{F}^{ab} = \frac{1}{2} \varepsilon^{abcd} F_{cd}$. Extension of this analysis to a nonabelian vector is straightforward, and as in the abelian case there is only one operator possible for each group at dim 5 level. Note that operator (1) is odd under *CPT* and even under charge conjugation. The equation of motion becomes

$$\square A_a = \frac{\xi}{M_{\text{Pl}}} \varepsilon_{abcd} n^b (n \cdot \partial)^2 F^{cd} \quad (2)$$

again with the gauge choice $\partial \cdot A = 0$. Further the right hand side has been reduced through ample use of the Bianchi identity. To identify the effect of the new term on the dispersion relation, we go to momentum space and select photons moving along the z axis with $k^a = (E, 0, 0, p)$. Then for transverse polarizations along the x and y axes,

$$\left(E^2 - p^2 \pm \frac{2\xi}{M_{\text{Pl}}} p^3 \right) (\epsilon_x \pm i\epsilon_y) \simeq 0 \quad (3)$$

where we have used $E \simeq p$ to leading order and chosen the “rest frame” where $n^a = (1, 0, 0, 0)$. Hence the sign of the cubic term is determined by the chirality (or circular polarization) of the photons. This leads to the rotation of the plane of polarization for linearly polarized photons, which may be used to bound ξ [3, 6]. Note that \mathcal{L}_γ is unique and hence the common approach[2, 4] of postulating a cubic dispersion relation which is chirality independent is incompatible with effective field theory.

Spinor: Consider a Dirac spinor for which the leading kinetic term is: $\mathcal{L}_0 = \bar{\Psi}(i\partial - m)\Psi$. The leading order equation of motion is just the Dirac equation: $(\partial + im)\Psi = 0$. In momentum space with $\Psi \sim \exp(-ik \cdot x)$, $(\not{k} - m)\Psi(k) = 0$. To modify the dispersion relation at $O(E^3)$, we consider new terms satisfying the constraints listed above. In this case, there are only two terms with the desired form

$$\mathcal{L}_f = \frac{i}{M_{\text{Pl}}} \bar{\Psi} (\eta_1 \not{n} + \eta_2 \gamma_5 \not{n}) (n \cdot \partial)^2 \Psi \quad (4)$$

Both operators break CPT, with η_1 being charge conjugation odd and η_2 charge conjugation even. After applying $(i\partial + m)$, the equation of motion takes the form

$$(\square + m^2)\Psi = \frac{i}{M_{\text{Pl}}} (\eta_1 + \eta_2 \gamma_5) (n \cdot \partial)^3 \Psi. \quad (5)$$

Here we have again dropped terms of order m/M_{Pl} or which vanish by the leading order equations. Hence the modified dispersion relation becomes

$$\left(E^2 - |\vec{p}|^2 - m^2 + \frac{|\vec{p}|^3}{M_{\text{Pl}}}(\eta_1 + \eta_2 \gamma_5)\right) \Psi = 0 \quad (6)$$

where we have used $E \simeq |\vec{p}|$ for high energies. At high energies (*i.e.*, $E^2 \gg m^2$), we can choose spinors as eigenspinors of the chirality operator and redefine coupling constants as $\eta_{L,R} = \eta_1 \mp \eta_2$. To introduce these operators for Standard Model, the chiral choice for η couplings would be required by gauge invariance. Previous studies[4, 5] considered only chirality independent dispersion relations for fermions and so implicitly fix $\eta_2 = 0$.

Above we have identified interesting operators which modify the dispersion relations at cubic order for vectors or fermions. The external tensor appearing in all of these operators takes the form $n^a n^b n^c$. For technical reasons to be addressed in the discussion section, we must replace this coupling by the traceless symmetric tensor $C^{abc} = n^a n^b n^c - \frac{1}{6}(n^a g^{bc} + \text{cyclic})$ in the following section. Note that this change implicitly introduces extra terms which do not satisfy all of the constraints listed above, *i.e.*, they are reducible by the leading order equations of motion. However, this replacement does not affect the dispersion relations in the regime $E \gg m$. One could also consider frame-dependent modifications of the interaction terms between, *e.g.*, photons and electrons. If we limit ourselves to the traceless symmetric coupling, C^{abc} , introduced here, there are in fact no additional interaction terms with dimension 5 beyond those implied by extending Eq. (4) with gauge-covariant derivatives.

3 Experimental constraints on dimension 5 operators

Evidence of modified dispersion relations for stable particles such as electrons, light quarks, and photons can be searched for using the astrophysical probes[2, 3, 4, 5, 6]. We begin here, however, by showing that impressive constraints can be imposed by considering terrestrial experiments. These indirect limits exploit the idea that the external four-vector n^a introduces a preferred frame that can not coincide with the laboratory frame on the Earth[13]. While the component n^0 is still dominant, the motion of the galaxy, Solar system and Earth will create spatial components $n_i \sim 10^{-3}$ for a terrestrial observer. Hence clock comparison experiments[14] or searches for spatial anisotropy[15] can impose stringent bounds on violations of Lorentz symmetry in this context. This approach was recently used to constrain the dispersion relation for nucleons[8]. Our constraints[9] apply to the fundamental fields of the Standard Model rather than presumably should have more direct connection to the Planck scale physics than nucleons.

Limits on the operators involving electrons and electron neutrinos are especially easy to derive. We use the fact that best tests[15] of directional sensitivity in the precession of electrons limit the size of interaction between the external direction and the electron spin at the level of 10^{-28} GeV. This immediately translates into the following limit on the coefficients η_L and η_R that parametrize the effective interaction of the form (4) for left-handed leptons and right-handed electrons[9]:

$$|\eta_L^e - \eta_R^e| \lesssim \frac{10^{-28} \text{ GeV } M_{\text{Pl}}}{m_e^2 |n_i|} \simeq 4, \quad (7)$$

where $M_{\text{Pl}} \equiv 10^{19}$ GeV. The combination, $\eta_L^e + \eta_R^e$, would be very weakly constrained here, as it does not appear in the electron spin Hamiltonian.

The absence of a preferred direction is checked with even greater precision using nuclear spin, which translates into more stringent limits on new operators for the light quarks. The photon operator (1) will also contribute because of the electromagnetic interactions inside the nucleon. To use the best experimental limits of 10^{-31} GeV on the coupling of n_i to neutron spins[14], we must relate the photon and quark operators with nucleon spin. First, let us introduce dimension 5 operators for the first generation of quarks, the left-handed doublet ψ_Q and right-handed singlets ψ_u and ψ_d :

$$\mathcal{L}_q = \frac{C^{abc}}{M_{\text{Pl}}} \sum_{i=Q,u,d} \eta_i \bar{\psi}_i \gamma_a \partial_b \partial_c \psi_i. \quad (8)$$

At the nucleon level, we estimate using standard QCD sum rules

$$\begin{aligned} \eta_{1,N} &= a_u(\eta_u + \eta_Q) + a_d(\eta_d + \eta_Q) \\ \eta_{2,N} &= b_u(\eta_u - \eta_Q) + b_d(\eta_d - \eta_Q) + b_\gamma \xi, \end{aligned} \quad (9)$$

where $\eta_{1(2),N}$ are the η_1 and η_2 couplings for nucleons defined in (4). Note that ξ enters only in the η_2 coupling for nucleons because both are even under the charge conjugation. In (9), $a_{u,d}$ and $b_{u,d}$ are the matrix elements that could be obtained as the moments of the experimentally measured structure functions[16]: $a_d \sim 0.4$, $a_u \sim 0.1$, $b_d \sim 0.1$, $b_u \sim -0.05$ for the neutron and charge inverted values for the proton. To relate the photon operator with nucleon, we use the simplest vector dominance model and obtain at one-loop level $b_\gamma \sim 0.13\alpha/(4\pi)$ for neutron and $b_\gamma \sim 0.24\alpha/(4\pi)$ for proton. Combining these results produces the following limit:

$$|(\eta_d - \eta_Q) - 0.5(\eta_u - \eta_Q) + 10^{-3}\xi| \lesssim 10^{-8} \quad (10)$$

Barring accidental cancellations, we can place separate limits on $\eta_{u,d} - \eta_Q$ at 10^{-8} and on ξ at 10^{-5} level. The orthogonal combinations $\eta_{u,d} + \eta_Q$ are less constrained because they enter only in the quadrupole coupling between the nuclear spin and external direction[8], and thus are suppressed by an additional factor of $|n_i| \sim 10^{-3}$.

So far we have neglected the fact that the low energy values for the couplings η_i and ξ taken at the normalization scale of 1 GeV do not coincide with the high-energy values for the same couplings generated at M_{Pl} . With a simple one-loop analysis of the renormalization group equations, we find that several bounds can be strengthened. Leaving the details for elsewhere[17], our results are: $|\eta_{Q,u,d}|, |\xi| \lesssim 10^{-6}$ and $|\eta_{L,R}^e| \lesssim 10^{-5}$. The constraints may also be improved by assuming degeneracies appropriate for grand unification at the GUT scale. On the experimental side, the recent progress[18] in the high-sensitivity atomic magnetometers that are unaffected by spin-exchange relaxation may lead to an improvement of terrestrial bounds on Lorentz violation by several orders of magnitude.

In summary, we have shown that effective field theory provides a framework where one can derive stringent bounds on Planck scale interactions from terrestrial experiments. The resulting limits enhance and generalize the terrestrial bounds obtained previously[8]. They are generally

far more sensitive than those previously derived by considering astrophysical phenomena[4] where the typical sensitivity is $O(1)$ in units of M_{Pl}^{-1} . A notable exception is the bound[3] $|\xi| \lesssim 2 \times 10^{-4}$ derived from the birefringence induced by \mathcal{L}_γ . Our bound[9] improved on this result by roughly an order of magnitude. However, with new astronomical data, this bound was recently updated[6] in a striking way to $|\xi| \lesssim 10^{-14}$.

Our limit (7) is already comparable to previous constraints from the astrophysical searches of vacuum photon decay and the absence of the vacuum Cerenkov radiation[4]. Further, the latter analysis did not consider polarization effects (*i.e.*, assumed $\eta_2 = 0$) and so this result provides a complementary constraint. Certainly our bounds inferred from the renormalization group analysis[9, 17] represent a major improvement in constraining these couplings. However, these bounds were already tightened to an extraordinary level by other recent work[5]. There, the observation of synchrotron radiation from the Crab nebula was used to infer $\eta_L^e + \eta_R^e \gtrsim -10^{-7}$. We should also note that combining various astrophysical constraints may yield $O(10^{-2})$ constraints on the difference $\eta_R^e - \eta_L^e$ subject to certain assumptions[6].

4 Discussion

The procedure of classifying operators according to their dimension proves very useful in analyzing sensitivity of various experiments to Lorentz violation. In such an analysis, one should always start from the *lowest* possible operators that may be induced by the short-distance physics. In the case of Lorentz violation, this expansion starts from dimension 3 [1], and hence we arrive at a certain “naturalness” problem for the whole approach: Why should the size of dimension 3 operators not be M_{Pl} , the scale of UV physics responsible for the breaking, rather than less than $\sim 10^{-30}$ GeV as various experimental tests of Lorentz invariance demand? Similarly, the naive expectation would be that Lorentz violation should appear with $O(1)$ strength at dimension 4, again in contradiction with experimental constraints. Perhaps, the full answer to this puzzle requires a proper understanding of the full dynamical picture of Lorentz violation emerging from high-energy scales. Without that we have to resort to a number of plausible explanations, none of them very compelling at present.

One possibility is that the properties of the theory at the UV scale are inconsistent with the presence of dimension 3 and 4 Lorentz-violating operators. For example, it can be easily argued that the exact supersymmetry at UV scale may completely forbid or at the very least severely restrict the possibility of dimension 3 and 4 operators. A classification of all supersymmetric Lorentz non-invariant operators is an interesting study on its own, but goes outside of the present discussion[19].

Another generic possibility to entertain in this approach is that the properties of the Lorentz-violating background itself lead to the suppression/exclusion of dimension 3 and 4 operators. For example, one might consider Lorentz violations as realized by a symmetric traceless rank-three tensor, C^{abc} , as was introduced at the end of section 2. In euclidean space, such tensor would correspond to an intrinsic *octupole* deformation of the vacuum. One finds that there are simply no dimension 3 or 4 operators which couple to this tensor. Beyond an absence of bare couplings, this implies that in a leading order calculation linear in C^{abc} can not generate

dimension 3 or 4 operators at any number of loops[17]. This was of particular importance in section 3 where we found the dimension 5 operators evolved logarithmically and while dimension 3 operators might have been expected to appear with $\Lambda_{UV}^2/M_{\text{Pl}}$ coefficients, they were in fact absent. Setting $C^{abc} = n^a n^b n^c - \frac{1}{6}(n^a g^{bc} + \text{cyclic})$ was a simplifying ansatz in our analysis and is not required by any underlying principles. Notice that if the dimension 5 operators coupled to $n^a n^b n^c$ rather than C^{abc} , one-loop graphs already generate, *e.g.*, $\Lambda_{UV}^2/M_{\text{Pl}} \bar{\psi} \not{n} \gamma_5 \psi$. The appearance of such quadratic divergence results because the new coupling is not in an irreducible representation of the Lorentz group and, as well as C^{abc} , contains a vector component which can couple to dimension 3 and 4 operators.

One can also consider higher order quantum corrections beyond linear order in C^{abc}/M_{Pl} . A priori, there is no reason to believe that this effect should be small. For example, a triple product $C_{ab}^c C^{abd} C_{cd}^e$ will transform as a vector, and therefore may generate dimension 3 operators with a coefficient $\Lambda_{UV}^4/M_{\text{Pl}}^{-3}$. If the scale of the ultraviolet cutoff in the loops is itself of the order of the Planck scale, then the resulting size of dimension 3 operator would be too large[17, 20]. Therefore, one is forced to assume that an effective cutoff in the loops with $\Lambda_{UV} \ll M_{\text{Pl}}$. If we require these induced dimension 3 operators to be smaller than dimension 5 operators at energy scales of 1 GeV, we must assume that $\Lambda_{UV} \lesssim \sqrt{M_{\text{Pl}}} \times 1\text{GeV} \sim 10^9\text{GeV}$. This scale is far larger than the electroweak scale and/or the scale of the supersymmetry breaking. Even though a generic investigation of the supersymmetric stabilization of the ultraviolet divergencies in the presence of the Lorentz violating operators is presently lacking, it certainly appears as a plausible scenario[19]. To summarize, one must regard both the appearance of C^{abc} and a low scale cutoff in the loops as essential ingredients in suppressing the appearance of Lorentz violations in lower dimension operators from quantum effects.

Following Ref. [8], it is amusing to compare the experimental constraints to semi-classical calculations which have appeared in the loop quantum gravity literature[21, 22, 23]. For a Dirac fermion, these studies[22] suggest that $\eta_1 = 0$ while η_2 is nonvanishing. However, the latter is suppressed by factors of M_L/M_{Pl} where $M_L \ll M_{\text{Pl}}$ is the coherence scale of the gravitational wave function. Therefore these calculations seem to be in agreement with the stringent experimental bounds[5, 6, 9] imposed on the fermion operators. In contrast, a separate analysis[21] suggests that ξ should be $O(1)$, which stands in stark contradiction with the bounds discussed above[3, 6, 9]. Hence naively the experimental bounds seem to be in conflict with the predictions of loop quantum gravity. However, it must be noted that the latter results are at present somewhat heuristic and so this apparent contradiction should not be taken too seriously.

Further recent calculations[24] suggest that the effects of quantum gravity may deform the Lorentz symmetry, *i.e.*, departures from the “standard” realization of Lorentz symmetries may be induced by Planck scale effects. This scenario, commonly known as “doubly special relativity” (DSR)[11], would seem a drastic departure from Lorentz breaking, in which our discussion was phrased. However, we will argue that many of the results are still applicable in this new scenario.

For simplicity, focus on photons and electrons which are central to many of the discussions of experimental bounds. First one must assume that the present preliminary investigations of DSR can be extended to provide something like a quantum field theory of photons and electrons

incorporating DSR. Now the first simple observation is that in the limit where $M_{\text{Pl}} \rightarrow \infty$, this DSR theory must reduce to ordinary QED. Then for finite M_{Pl} , the low-energy or long-wavelength physics can still be described using standard techniques of effective field theory. Hence if the dispersion relations relevant for photons and electrons are modified at $O(p^3)$ (and we assume that rotational invariance is unmodified), then it must be by the appearance of dimension 5 interactions of the form given in Eqs. (1) and (4). In this context, one should not think that these new interactions break Lorentz invariance. Rather the apparent Lorentz violations induced by standard Lorentz transformations on these terms must be compensated by corrections to these transformations when they act on the standard kinetic terms. In fact, the latter observation may be used as a strategy to infer the leading corrections to the Lorentz transformations[25]. In any event, the discussion of section 2 is equally applicable to DSR scenarios as to scenarios of Lorentz breaking.

One must, of course, be more cautious when considering the experimental bounds quoted above, as many of these rely on the existence of a preferred frame[10]. However, we argued above that Eq. (1) still gives the leading modification of the photon action in a DSR scenario, and hence the $O(p^3)$ modification of the dispersion relation is photon-chirality dependent.¹ Further, the astronomical tests involving the resulting birefringence for photons provide purely kinematical constraints on the theory. That is, one relies on the magnification of small but unusual effects in the propagation of photons as they cross cosmological distances, without reference to a preferred frame. Hence these bounds[3, 6] are again applicable to DSR scenarios. Given the present experimental bound[6] that $|\xi| \lesssim 10^{-14}$, it seems that cubic modifications of the photon dispersion relation are essentially ruled out for DSR constructions.

Another important aspect of the effective field theory approach to Lorentz violation is the correct identification of the infrared degrees of freedom and the possibility of generalizing this framework to general relativity[12]. Attempts to generalize Lorentz breaking operators to general relativity would necessarily supply C^{abc} with coordinate dependence and a corresponding kinetic term, which in turn would contribute to the energy density. There are several arguments supporting the idea that Lorentz violating tensors must be accompanied by massless (or nearly massless) particles in the spectrum. Many explicit examples involving Lorentz violation point towards this result, including noncommutative backgrounds in string theory[26], super-light axion backgrounds[27] and ghost condensation[28]. The exchange by new light degrees of freedom may lead to the appearance of a new coherent interaction, *i.e.*, a “fifth force”. Therefore, quite separate bounds on Lorentz violation may emerge from gravitational physics and cosmology. Elucidating this question is also important for interpreting the bounds on Lorentz violation that come from cosmological distances, over which a “constant” background would most certainly be changing.

In conclusion, we would like to stress again that there remains much work to be done both on the theoretical description from quantum gravity and the phenomenological constraints. However, it is truly remarkable that present-day precision experiments and astrophysical observations can already confront quantum gravity calculations with concrete and stringent observational bounds.

¹It is curious that certain arguments seem to imply this would be inconsistent with DSR[10].

Acknowledgments

RCM would like to thank the organizers of QTS3 for the opportunity to present this material. We would like to thank Giovanni Amelino-Camelia, Cliff Burgess, Laurent Freidel, Stefan Groot Nibbelink, Ted Jacobson, Jerzy Kowalski-Glikman, Joe Lykken, Seth Major, Guy Moore, Ann Nelson, Adam Ritz, Mike Romalis, Subir Sarkar and Lee Smolin for many useful comments and conversations over the course of this work. This research is supported in part by NSERC of Canada, Fonds FCAR du Québec and PPARC UK.

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